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Can MSSM Particle be the Inflaton?

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Abstract

We consider the possibility of using one of the D -flat directions in the minimal supersymmetric standard model (MSSM) as the inflaton. We show that the flat direction consisting of (first generation) left- and right-handed up-squarks as well as the up-type Higgs boson may play the role of the inflaton if dominant part of the up-quark mass is radiatively generated from supersymmetric loop diagrams. We also point out that, if the R -parity violating Yukawa coupling is of $O(10^{-7})$, R -odd D -flat directions may be another possible candidate of the inflaton. Such inflation models using D -flat directions in the MSSM are not only testable with collider experiments but also advantageous to resolve the problem how the inflaton reheats the universe.

Inflation [1] is now one of the most important ideas in cosmology. Inflation not only solves the serious horizon and flatness problems but also provides a viable scenario of generating the origin of cosmic density fluctuations.^{#1} In particular, precise measurement of the anisotropy of the cosmic microwave background (CMB) by the Wilkinson Microwave Anisotropy Probe (WMAP) suggests that the primordial density fluctuations are almost purely adiabatic and scale-invariant [3], which are predictions of (some classes of) inflationary models. In the inflationary models, a scalar field, called “inflaton,” is introduced to realize the inflationary epoch of the universe. During inflation, potential energy of the inflaton gives the dominant part of the energy density. In order to realize the quasi de Sitter universe with the potential energy of the inflaton, its kinetic energy should be much smaller than the potential energy during inflation. Consequently, we are led to the paradigm of slow-roll inflation, where inflation is driven by the potential energy of slowly evolving scalar field.

It is non-trivial to find a scalar field which satisfies the slow-roll condition, and from the particle-physics point of view, it is important to find good candidates of the inflaton. In particular, it is interesting to ask if the inflaton can be observed at high-energy collider experiments in any scenario of inflation. In many classes of models, however, a new scalar field is introduced as the inflaton. Usually, such scalar fields do not belong to the standard model and have very weak interactions with the standard-model particles, which makes it very difficult to find and study the inflaton with collider experiments. Moreover, since the interactions between such inflaton and standard-model particles cannot be determined, it must be given *ad hoc* by hand, which obscures the thermal history of the universe. Thus the reheating temperature, for instance, cannot surely be estimated. In order to construct a testable and economical model of inflation, it is desirable to find a candidate of the inflaton in the scalar fields which are in some sense familiar to us. In fact, in the standard model, the only scalar field is the Higgs boson. As we will see later, however, it is known that the Higgs boson in the standard model cannot be the inflaton since its (quartic) coupling is too large to generate cosmic density fluctuations that are consistent with the observations. In addition, in the standard model, the potential of the Higgs boson is significantly affected by radiative corrections and hence we cannot expect flat enough potential which is crucial to cause inflation long enough. Thus, we should conclude that it is impossible to find a viable candidate of the inflaton in the standard-model fields. In fact, the second point, effects of radiative corrections to the inflaton potential, is in general a serious problem in constructing inflation models.

In order to control the radiative corrections, it is often the case that inflation models are considered in supersymmetric framework; indeed, in supersymmetric models, quadratic divergences cancel out between bosonic and fermionic loops and the flatness of the potential can be guaranteed. Thus, in considering the inflation, supersymmetry is likely to play very important roles, and in this letter, we adopt (low-energy) supersymmetry. If we super-

^{#1}Another possibility of generating cosmic density fluctuations may be to adopt the “curvaton” scenario [2]. Here, we do not consider such a possibility and assume that the cosmic density fluctuations are totally generated from the primordial fluctuation of the inflaton.

symmetrize the standard model, various scalar particles are introduced as superpartners of quarks and leptons. Those scalars may play the role of inflaton. It is worth noting that, if this is the case, the reheating processes into the standard-model particles are obvious, which makes it possible to study both the inflationary and thermal history of the universe only with the low-energy “known” physics. Actually there was an early attempt to build inflation models along this line [4]. However, as discussed later, it contained some difficulty in producing the right amount of density fluctuations, so the authors of Ref. [4] had to introduce additional mini-inflation. Here we would like to pursue another possibility.

In this letter, we consider the possibility of using scalar fields in the minimal supersymmetric standard model (MSSM) as the inflaton (which is denoted as ϕ hereafter). In particular, we will discuss that the D -flat direction consisting of first-generation left- and right-handed up-type squarks as well as the up-type Higgs boson may play the role of the inflaton if the up-quark mass is radiatively generated. We will also see that, if the R -parity violating Yukawa coupling is of $O(10^{-6} - 10^{-7})$, R -parity violating D -flat directions may also be the inflaton.

We start with discussing possible scenario of inflation within the MSSM. In the framework of the slow-roll inflation, the amplitude of the inflaton field varies during and after the inflation. The change of the inflaton amplitude is typically of $O(M_*)$, where $M_* \simeq 2.4 \times 10^{18}$ GeV is the reduced Planck scale.^{#2} Thus, inflaton originating from the MSSM sector should have an amplitude of $O(M_*)$ during inflation. Consequently, we are forced to consider the chaotic inflation [5] within the MSSM assuming that the MSSM field corresponding to the inflaton has an amplitude of the order of the Planck scale during inflation. Later we will discuss how to realize such a large field value.

Now, we discuss the observational constraints on the inflaton potential. In the simplest scenario, the cosmic density fluctuations are parameterized by the curvature perturbation \mathcal{R} which depends on the inflaton potential V as

$$\mathcal{R}(k) = \left[\frac{H_{\text{inf}}}{2\pi} \frac{3H_{\text{inf}}^2}{V'} \right]_{k=aH_{\text{inf}}}, \quad (1)$$

where H_{inf} is the expansion rate H of the universe during inflation, which is related to the potential energy of the inflaton during inflation as $H_{\text{inf}}^2 = V/3M_*^2$, while V' is the derivative of the inflaton potential with respect to the inflaton field. Here, notice that this quantity is evaluated at the time when the fluctuation exits the horizon during inflation; k and a denote the wave-number (for the comoving coordinate) and scale factor, respectively.

From the measurements of the CMB anisotropies (as well as from other observations), we can obtain constraints on $\mathcal{R}(k)$. The most important constraint is on its normalization. Due to the fact that $\Delta T/T \sim O(10^{-5})$, the typical size of \mathcal{R} is also constrained to be $O(10^{-5})$. More precisely, the WMAP team [3] estimated it as $|\mathcal{R}(k)|^2 = 2.95 \times 10^{-9} A$

^{#2}We consider neither the so-called small-field models (for example, new inflation) nor hybrid models within the MSSM. Generally speaking, both require the vacuum to be such that the standard-model gauge symmetries are spontaneously broken, which is hardly justified. Here and hereafter we concentrate on the so-called large-field models (*i.e.*, chaotic inflation).

with $A = 0.9 \pm 0.1$ at $k = 0.05 \text{ Mpc}^{-1}$, assuming power-law ΛCDM model. If the inflaton potential has the parabolic form $V = \frac{1}{2}M_\phi^2\phi^2$, the inflaton mass is required to be $M_\phi \sim O(10^{13} \text{ GeV})$. Obviously, such a heavy scalar field does not exist in the MSSM.

We consider the next possibility where the inflaton potential is quartic:

$$V = \frac{1}{4}\lambda\phi^4, \quad (2)$$

where λ is a dimensionless coupling constant. Then, the normalization of the primordial density fluctuation requires $\lambda \sim O(10^{-13})$ as follows. The value of the inflaton field is related to the e -folding number N as $\phi \simeq \sqrt{8N}M_*$. (See Eq. (14) below). Thus the curvature perturbation at the horizon exit is evaluated as $\mathcal{R} \simeq 0.3\sqrt{\lambda}N^{3/2}$. Equating this with the WMAP result, we obtain $\lambda \sim 10^{-13}$, where $N \sim 60$ is used. Before studying the MSSM cases, we should comment on the non-supersymmetric case; from this constraint, it is obvious that the Higgs boson cannot play the role of the inflaton since, if the quartic coupling of the Higgs boson is as small as $O(10^{-13})$, Higgs mass becomes $O(10^{-4} \text{ GeV})$ which is unacceptably smaller than the present experimental bound.

In the supersymmetric case, such a small coupling for quartic interaction cannot be realized if the potential is lifted by the gauge D -term interactions, since, if so, the coupling constant λ becomes of the order $O(g^2)$ where g is the gauge coupling constant in the standard model. Therefore we focus our attention on the D -flat directions. For D -flat directions, we have to be more careful since behaviors of the potential depend on which flat direction we consider. In the MSSM, Yukawa interactions exist in the superpotential to generate the fermion masses. Such Yukawa interactions lift some of the D -flat directions. In addition, we can also find several D -flat directions which are not affected by the Yukawa interactions associated with the fermion masses; without R -parity violation, such D -flat directions are only lifted by the effects of supersymmetry breaking.^{#3} (See Ref. [6] for the details.)

We first consider the D -flat direction lifted by the R -parity conserving Yukawa interactions. D -flat directions are parameterized by gauge invariant monomial of the superfields. We denote the MSSM superfields as $Q(\mathbf{3}, \mathbf{2}, \frac{1}{6})$, $U(\mathbf{3}^*, \mathbf{1}, -\frac{2}{3})$, $D(\mathbf{3}^*, \mathbf{1}, \frac{1}{3})$, $L(\mathbf{1}, \mathbf{2}, -\frac{1}{2})$, $E(\mathbf{1}, \mathbf{1}, 1)$, $H_u(\mathbf{1}, \mathbf{2}, \frac{1}{2})$, and $H_d(\mathbf{1}, \mathbf{2}, -\frac{1}{2})$, where we show the quantum numbers for the $SU(3)_C \times SU(2)_L \times U(1)_Y$ gauge group in the parentheses. Then, the relevant part of the MSSM superpotential is given by

$$W = [Y_U]_{ij}Q_iU_jH_u + [Y_D]_{ij}Q_iD_jH_d + [Y_E]_{ij}L_iE_jH_d, \quad (3)$$

where i and j are generation indices. If we consider the D -flat directions represented as $Q_iU_jH_u$, $Q_iD_jH_d$, and $L_iE_jH_d$, those D -flat directions acquire quartic potential due to the Yukawa interactions, as given in Eq. (2). In order to relate the Yukawa coupling constants to λ in Eq. (2), we express those D -flat directions with a complex scalar field

^{#3}Here, we assume that coefficients of non-renormalizable terms are suppressed enough to be neglected. This may be explained by the R -symmetry, assigning R -charge $\frac{2}{3}$ to each MSSM chiral superfields.

Φ , for example,

$$Q_i = \frac{1}{\sqrt{3}} \begin{pmatrix} \Phi \\ 0 \end{pmatrix}, \quad U_j = \frac{1}{\sqrt{3}} \Phi, \quad H_u = \frac{1}{\sqrt{3}} \begin{pmatrix} 0 \\ \Phi \end{pmatrix}. \quad (4)$$

Self quartic coupling constants of those flat directions, denoted as $\lambda_{Q_i U_j H_u}$, $\lambda_{Q_i D_j H_d}$, and $\lambda_{L_i E_j H_d}$, respectively, are then given by

$$\lambda_{Q_i U_j H_u} = \frac{1}{3} |[Y_U]_{ij}|^2, \quad \lambda_{Q_i D_j H_d} = \frac{1}{3} |[Y_D]_{ij}|^2, \quad \lambda_{L_i E_j H_d} = \frac{1}{3} |[Y_E]_{ij}|^2, \quad (5)$$

where we defined $\phi \equiv \sqrt{2} \text{Re } \Phi$. Thus, if one of these coupling constants is of $O(10^{-13})$, we may have inflaton candidates within the MSSM particles.

If the fermion masses totally originate from the superpotential given in Eq. (3), Yukawa coupling constants are evaluated as

$$[\hat{Y}_U]_{ii} = \frac{m_{u_i}}{v \sin \beta}, \quad [\hat{Y}_D]_{ii} = \frac{m_{d_i}}{v \cos \beta}, \quad [\hat{Y}_E]_{ii} = \frac{m_{e_i}}{v \cos \beta}, \quad (6)$$

where \hat{Y}_U , \hat{Y}_D , and \hat{Y}_E are Yukawa matrices in the diagonalized basis, while m_{u_i} , m_{d_i} , and m_{e_i} are fermion masses for corresponding Yukawa couplings. In addition, $v \equiv \sqrt{\langle H_u^2 \rangle + \langle H_d^2 \rangle} \simeq 174 \text{ GeV}$ while $\tan \beta \equiv \langle H_u \rangle / \langle H_d \rangle$. Yukawa coupling constants for the second and third generation quarks and leptons are much larger than $\sim 10^{-6}$, and hence we consider the possibility of using the first generation squarks and/or sleptons as the inflaton. (Hereafter, we consider only the first generation squarks and sleptons, and drop the generation indices for simplicity unless otherwise mentioned.) If we estimate the Yukawa coupling constants for the up and down-quarks as well as electron using Eq. (6), we obtain

$$y_u \simeq \frac{8.6 \times 10^{-6}}{\sin \beta} \times \left(\frac{m_u}{1.5 \text{ MeV}} \right), \quad y_d \simeq \frac{2.9 \times 10^{-5}}{\cos \beta} \times \left(\frac{m_d}{5 \text{ MeV}} \right), \quad y_e \simeq \frac{2.9 \times 10^{-6}}{\cos \beta}. \quad (7)$$

In the MSSM, it is often the case that, in order to evade the Higgs-mass constraint, relatively large value of $\tan \beta$ is required [7]. Then, y_d and y_e are likely to be larger than 10^{-6} . In fact, if we adopt the up-quark mass $m_u = 1.5 - 4.5 \text{ MeV}$ [8],^{#4} we obtain $y_u \geq 8.6 \times 10^{-6}$, leading to $\lambda_{QUH_u} \geq 2.5 \times 10^{-11}$, and hence even the up-squark cannot play the role of the inflaton.^{#5} One might well give up the simple $\lambda\phi^4$ model and add another mini-inflation to reconcile the predicted magnitude of density fluctuations with the observed one [4]. However, here we would like to stick to the $\lambda\phi^4$ model, since it predicts the density fluctuations with very small uncertainty as shown later.

^{#4}Here, we do not consider renormalization group running which suppresses the fermion masses at higher energy scale. Even if the renormalization group effects are taken into account, our discussions are qualitatively unchanged.

^{#5}The quark mass ratios are constrained by imposing a limit on next-to-leading order corrections in the chiral perturbation theory. If the corrections become sizable, the up quark mass can be much smaller than the lower limit $\sim 1.5 \text{ MeV}$ [9]. Then, y_u can be as small as $\sim O(10^{-7})$. In this case, we do not have to consider the radiatively induced up-quark mass as we will discuss in the following.

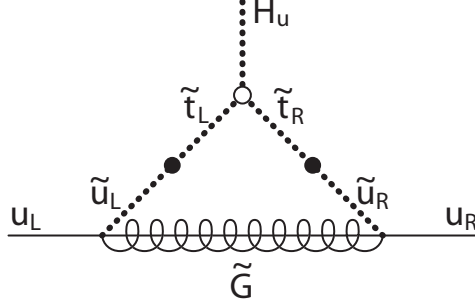


Figure 1: Loop diagram contributing to the up-quark mass in the MSSM. Black dots represent the insertion of the off-diagonal elements of the squark mass matrices while the vertex with the open circle is from the insertion of the trilinear coupling A_t .

So far, we have assumed that the Yukawa interactions in the superpotential are the only sources of the fermion masses. In the MSSM, however, radiative corrections due to the supersymmetric loop also affect the fermion masses. If such an effect is large enough to explain the dominant part of the fermion mass, Yukawa coupling constant in the superpotential can be smaller than the naive expectation. Indeed, for the up-quark mass, contribution from the squark-gluino loop diagram may become sizable with non-vanishing off-diagonal elements of the left- and right-handed squark mass matrices, as given in Fig. 1. Relevant part of the supersymmetry breaking terms contributing to this diagram is given by

$$\mathcal{L} = \Delta m_{\tilde{u}_L \tilde{t}_L}^2 \tilde{u}_L \tilde{t}_L^* + \Delta m_{\tilde{u}_R \tilde{t}_R}^2 \tilde{u}_R \tilde{t}_R^* + A_t \tilde{t}_L \tilde{t}_R H_u + \frac{1}{2} m_{\tilde{G}} \tilde{G} \tilde{G} + \text{h.c.}, \quad (8)$$

where $\tilde{u}_{L,R}$ and $\tilde{t}_{L,R}$ are left- and right-handed up- and top-squarks, respectively, while \tilde{G} denotes the gluino. Hereafter, we approximate that the masses of the squarks are all degenerate for simplicity, and we define

$$x \equiv \frac{m_{\tilde{G}}^2}{m_{\tilde{q}}^2}, \quad \delta_{13}^{(L,R)} \equiv \frac{\Delta m_{\tilde{u}_{L,R} \tilde{t}_{L,R}}^2}{m_{\tilde{q}}^2}, \quad a_t \equiv \frac{A_t}{y_t m_{\tilde{q}}}, \quad (9)$$

where $m_{\tilde{q}}$ is the squark mass, and y_t is the top Yukawa coupling constant. Then, the loop contribution to the up-quark mass is given by

$$m_u^{(\text{loop})} = \frac{1}{36\pi^2} g_3^2 y_t a_t \delta_{13}^{(L)} \delta_{13}^{(R)} \langle H_u \rangle \frac{x^{1/2} (x^3 - 6x^2 + 3x + 2 + 3x \ln x)}{(x-1)^4}, \quad (10)$$

where g_3 is the $SU(3)_C$ gauge coupling constant.

Size of the off-diagonal elements of the squark mass matrices are constrained from flavor changing processes. In particular, for $\delta_{13}^{(L)}$, one finds $\delta_{13}^{(L)} < 4.6 \times 10^{-2}$ (9.8×10^{-2}),

2.3×10^{-1}) for $x = 0.3$ (1.0 and 4.0) and $m_{\tilde{q}} = 500$ GeV [10].^{#6} Constraint on $\delta_{13}^{(R)}$ is, on the other hand, not available since we do not have precise measurement of the flavor changing decay of the top quark. Thus, the combination $\delta_{13}^{(L)} \delta_{13}^{(R)}$ can be of $O(0.1)$ without conflicting experimental constraints.^{#7} Then, for example for the case where $m_{\tilde{G}} \simeq m_{\tilde{q}}$, we obtain

$$m_u^{(\text{loop})} \simeq \frac{1}{72\pi^2} g_3^2 y_t a_t \delta_{13}^{(L)} \delta_{13}^{(R)} \langle H_u \rangle \simeq 3.6 \text{ MeV} \times y_t a_t \sin \beta \left(\frac{\delta_{13}^{(L)} \delta_{13}^{(R)}}{10^{-2}} \right). \quad (11)$$

Thus, the up-quark mass of the size of a few MeV can be generated from the loop effect with reasonable value of $a_t \sim O(1)$. Then, the up-Yukawa coupling constant in the superpotential can be of $O(10^{-6} - 10^{-7})$. Notice that the loop-induced up-quark mass is not suppressed even if the masses of the superparticles become much larger than the electroweak scale (as far as all the supersymmetry breaking parameters are of the same order). Thus, if one wishes to push up $m_u^{(\text{loop})}$ without affecting flavor-changing processes, one possibility is to assume (relatively) large masses for the superparticles. We can also consider radiative correction to the down quark mass. In this case, A_t should be replaced by $y_b \mu_H$ in the calculation, where y_b and μ_H are bottom-quark Yukawa coupling and supersymmetric Higgs mass, respectively, and the generation mixings are provided by the down-type squark mass matrices. Thus, radiative correction to the down-quark mass may become sizable when $\tan \beta$ is large. Constraints on the off-diagonal elements of the down-type scalar quark mass matrices are, however, more stringent than those on the up-type scalar quark mass matrices. Thus, it is more difficult to radiatively generate the down-quark mass.

The above arguments open a window to use the D -flat direction QUH_u as the inflaton. Thus, hereafter, we concentrate on the evolution of the universe for the case where D -flat direction QUH_u plays the role of the inflaton, assuming that the dominant part of the up-quark mass is from the radiative correction.

The potential of QUH_u is now given by

$$V(\phi) = \frac{1}{2} m_\phi^2 \phi^2 + \frac{\lambda}{4} \phi^4, \quad (12)$$

where $m_\phi \sim O(100 \text{ GeV})$ is the soft supersymmetry breaking mass. These two terms become comparable at $\phi_m \equiv \sqrt{2/\lambda} m_\phi$. Since the mass term can be neglected for large value of $\phi \gg \phi_m$, its dynamics is just the same as the well-known $\lambda \phi^4$ model. First let

^{#6}Constraint on $\delta_{13}^{(L)}$ here is from the mass difference of the neutral B_d mesons. Such a constraint is obtained from the basis where the down-type squarks are diagonalized. Two basis are related by the Kobayashi-Maskawa matrix, but we neglect the difference of the two basis because the change of the constraint is small enough to be neglected.

^{#7}If the supersymmetry breaking parameters have large complex phases, too large neutron electric dipole moment may be induced in this model. Here, we consider the case where the phases in the supersymmetry breaking parameters are small enough. We also assume that the mixing parameters for the first and second generation squarks are suppressed in order to evade the constraints from the $K^0 - \bar{K}^0$ mixing.

us briefly review the $\lambda\phi^4$ model and discuss the observational constraint on that. In this model, the slow roll parameters are given by

$$\epsilon \equiv \frac{1}{2}M_*^2 \left(\frac{V'}{V} \right)^2 = 8 \frac{M_*^2}{\phi^2}, \quad \eta \equiv M_*^2 \frac{V''}{V} = 12 \frac{M_*^2}{\phi^2}, \quad (13)$$

where we substituted the quartic potential Eq. (2). When the inflaton ϕ becomes equal to $\phi_{\text{end}} \equiv 2\sqrt{3}M_*$, the slow-roll condition breaks down, leading to the end of inflationary epoch. The e -folding number N is given by

$$N \equiv \int_t^{t_{\text{end}}} H dt \simeq \int_{\phi_{\text{end}}}^{\phi} \frac{3H^2}{V'} d\phi = \frac{1}{8M_*^2} (\phi^2 - \phi_{\text{end}}^2), \quad (14)$$

where we used the slow-roll approximation in the second equality.

In order to impose observational constraints on inflation models, it is necessary to evaluate several parameters that characterize the density fluctuations generated during inflation. They are scalar spectral index n_s , its running $dn_s/d\ln k$, and tensor-to-scalar ratio r , apart from the normalization that determines the value of λ . These are related to the e -folding number as

$$n_s = 1 - \frac{3}{N + \frac{3}{2}}, \quad \frac{dn_s}{d\ln k} = -\frac{3}{(N + \frac{3}{2})^2}, \quad r = \frac{16}{N + \frac{3}{2}}. \quad (15)$$

Thus we must precisely evaluate N in order to make a comparison between predictions and observational results. Since the value of N necessary to solve the horizon and flatness problems depends on the thermal history of the universe, we must specify the reheating processes. Fortunately, however, it is possible to determine the e -folding number as $N \simeq 64$ with small uncertainty in the case of $\lambda\phi^4$ model. The reason is that the energy density of inflaton oscillation decreases just like radiation after the end of inflation, which applies only to the quartic model. The explicit expression for N is given by [11]

$$N = 62 - \ln \left(\frac{k}{a_0 H_0} \right) - \ln \left(\frac{10^{16} \text{ GeV}}{V_k^{1/4}} \right) + \ln \left(\frac{V_k^{1/4}}{V_{\text{end}}^{1/4}} \right), \quad (16)$$

where V represents the potential energy of the inflaton, and the subscripts “0,” “ k ,” and “end” are for variables at present, at the time when the fluctuation with the wave-number k exits the horizon, and at the end of the inflation, respectively. The right hand side of Eq. (16) gives $N \simeq 64$ for $k = a_0 H_0$, irrespective of the details of the reheating. Here we have assumed that the reheating processes complete before the amplitude of the inflaton becomes smaller than ϕ_m , which will be justified below. Substituting $N \simeq 64$ into Eqs. (15), we obtain $n_s \simeq 0.95$ and $r \simeq 0.24$ with negligible running of n_s .

Constraints on the inflationary models driven by a single slow-rolling scalar field from the recent observations including the WMAP have been studied extensively [12, 13, 14]. Although Ref. [12] claimed that $\lambda\phi^4$ model lies in the region marginally excluded by the

WMAP data in combination with smaller scale CMB and large scale structure survey data, more detailed analyses [13, 14] showed that the model cannot be excluded by the WMAP data alone for $N \gtrsim 40$. In addition, the recent systematic study using both the Sloan Digital Sky Survey and WMAP demonstrated that $\lambda\phi^4$ model is still allowed [15]. Thus there is no reason to disregard $\lambda\phi^4$ model at present, and one could well contend that the model is of much interest since it is on the edge.

The next discussion concerns the decay processes of the inflaton. After inflation ends, the inflaton oscillates around its origin until the decay completes. In the usual chaotic inflation models, the reheating process proceeds through nonperturbative particle creation (preheating) [16]. During the oscillation, the particles coupled to the inflaton are produced when the so-called adiabatic condition is violated, *i.e.*, $\dot{\omega}/\omega^2 \gtrsim 1$, where ω is the effective frequency of the produced particles. In our model, the inflaton field is complex, and its nontrivial trajectory in the potential may lessen the efficiency of the preheating process [17]. The possible source of the nontrivial orbit of the inflaton in our case is the A -term such as

$$V_A = a_u y_u m_{\tilde{q}} \Phi^3 + \text{h.c.}, \quad (17)$$

where a_u is a constant of $O(1)$. The effect of this A -term is, however, so small that the inflaton Φ exhibits almost straight-line motion on the complex plane. Therefore, when the inflaton comes closest to the origin, its amplitude is much smaller than the critical value, below which the adiabaticity condition is violated. Hence the preheating should proceed in the same way as a real scalar field. In fact we have confirmed numerically that the instability band almost coincides with that in the case of a real scalar field, even in the presence of the A -term. Also, since the preheating occurs very efficiently and ceases within several oscillations as shown below, cumulative disturbance of the homogeneous motion caused by the A -term can be neglected.

We would like to focus on the four-point scalar interaction with the stop among many interactions the inflaton feels. When the inflaton ϕ first reaches $\phi = 0$, the particle production occurs and the stops are generated, typically with the momentum $k_{\text{res}} \sim y_t^{1/2} \lambda^{1/4} \phi_0$ and the occupation number $n_k \sim O(1)$ for $k \sim k_{\text{res}}$, where ϕ_0 is the amplitude of the oscillation. Since the stop obtains an effective mass $m_{\tilde{t}} \sim y_t |\phi|$ through the four-point interaction, generated stops are fattened as the amplitude of inflaton increases, and become as heavy as $\sim y_t M_*$ around the endpoint of the oscillation of the inflaton. The stop can then decay into two fermions, wino and (left-handed) bottom quark, for example, because the time scale of the decay is much shorter than that of the oscillation of the inflaton. Similar process occurs when ϕ reaches $\phi = 0$ again. Each time ϕ passes its origin, stops are generated, and they decay when ϕ reaches around its maximal value. This type of preheating is known as “instant preheating” [18]. Applying the result of Ref. [18] to our case, the amount of the energy dissipated by each decay of the stops is comparable to that originally stored in the inflaton. Thus, the decay processes of the inflaton proceed very efficiently and complete within several oscillations.

After that, the decay products are quickly thermalized through decays, scatterings, and annihilations by gauge interactions, which proceed very efficiently. The reheating

temperature T_{RH} is expected to be very high: $T_{RH} \sim 0.1\lambda^{1/4}M_* \sim 10^{14}\text{GeV}$. Such high reheating temperature might lead to the overproduction of dangerous relics like gravitinos [19]. In order to avoid this problem, we assume that one of the followings is realized. One solution is to assume large late-time entropy production from, for example, thermal inflation [20]. Another is to have a relatively heavy gravitino mass, $m_{3/2} \simeq 10 - 100\text{TeV}$, so that the gravitinos can decay well before the big bang nucleosynthesis (BBN) epoch, $T_{\text{BBN}} \sim 1\text{MeV}$. However, this does not remedy the situation if the lightest supersymmetric particle (LSP) is the standard bino-like LSP, since those produced from the decay of gravitinos would overclose the universe. This difficulty can be evaded by the introduction of a supersymmetric partner with a mass much lighter than 100GeV . One possibility may be the axino, superpartner of the axion [21].

So far we have not mentioned how the flat direction can take the value of $O(M_*)$ or larger avoiding the Hubble-induced mass term that prevents the flat direction from slow-rolling (so called η -problem). In the minimal supergravity model, the scalar potential includes an exponential factor which essentially precludes any scalar amplitudes larger than M_* , and scalar masses of the order of H are generated. An idea to circumvent these obstacles to construct a successful chaotic inflation model is to introduce the Heisenberg symmetry [22] under which the inflaton ϕ and a chiral field z transform as follows,

$$\delta z = \epsilon^* \phi, \quad \delta \phi = \epsilon, \quad (18)$$

where ϵ is a complex parameter. We can construct an invariant combination y from z and ϕ as

$$y \equiv z + z^* - \phi^* \phi. \quad (19)$$

Imposing the Heisenberg symmetry in the Kähler potential, it is written only with y , *i.e.*, $K = f(y)$. It is easy to see that y and ϕ should be regarded as independent variables since the kinetic terms are diagonalized for these variables [23]. It was shown that this symmetry protects the flatness of the inflaton potential from both the exponential growth and Hubble-induced mass term. It should be emphasized that the introduction of a new degree of freedom y is the price for keeping the potential flat. Central to this issue is the problem how to stabilize y . The dynamics of y with a specific form of $f(y) = \frac{3}{8} \ln y + y^2$ was discussed in Ref. [24], and it was found that the value of y is fixed during the inflation. In addition, another mechanism of fixing y using the radiative corrections was proposed in Ref. [23] in the case of the no-scale supergravity model [25] where the Kähler potential takes a special form as $K = -3 \ln y$. In both cases, the potential of the inflaton is same as that in the global supersymmetry case as long as y is fixed. In this letter, we just assume y is somehow fixed and remains constant for successful chaotic inflation.^{#8}

^{#8}Notice that, even in this framework, soft supersymmetry breaking terms required for our mechanism can be generated by introducing a new supersymmetry breaking field x (with a slight modification of the Kähler potential). For example, let us consider the Kähler potential of the form $K = f(y) + a_{ij}|x|^2(\phi_i^* \phi_j + \text{h.c.})$ with a_{ij} being constants, ϕ_i MSSM chiral multiplets, and $y = z + z^* + \phi_i^* \phi_i + |x|^2$. This Kähler potential does not have the Heisenberg symmetry if $a_{ij} \neq 0$. Previous arguments, however, still hold since

Finally, we comment on the case with the R -parity violation. In the MSSM, there are D -flat directions which are not lifted by the Yukawa interactions given in Eq. (3). Such flat directions are parameterized by the monomial of the superfields with odd R -parity, and can be lifted if the R -parity is broken. Choosing relevant D -flat direction, R -parity breaking Yukawa coupling can be as large as $\sim O(10^{-6} - 10^{-7})$ without conflicting experimental bounds if the baryogenesis takes place after the sphaleron interaction becomes ineffective [26]. Thus, if we adopt R -parity violation of this size, such an R -parity violating D -flat direction can be another candidate of the inflaton. The decay of R -parity violating D -flat directions occurs in the same way as the previous case via the “instant preheating.” Since they do not necessarily have interactions with stops, the four-point scalar interaction with the largest coupling constant is expected to come from the D -term. In particular, there are four-point scalar interactions $\sim g^2 \chi^2 \phi^2$, where χ represents the field orthogonal to the flat direction.

In summary, we have investigated whether the inflation can be embedded in the MSSM sector, and found that the D -flat direction consisting of the first generation left- and right-handed up squarks and the up-type Higgs boson may be the inflaton if the up-quark mass predominantly comes from the one-loop threshold correction to the up-Yukawa coupling constant. The dynamics of the inflaton is almost the same as the $\lambda\phi^4$ model, which has attracted much attention recently. Since the inflaton in our model consists of the MSSM particles, it is not only minimal but also testable at high-energy collider experiments. In particular, if the QUH_u flat direction plays the role of the inflaton, mixings of the up- and top-squarks should be large. This is an interesting check point of our model and can be tested by collider experiments as well as precise measurements of flavor-changing processes.

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$|x| \ll M_*$ is realized during the inflation because of the Hubble-induced mass of x ; with $|x| \ll M_*$, inflaton potential does not change. In addition, if the F -component of x is non-vanishing (and large enough) in the vacuum, soft supersymmetry breaking scalar masses squared (including the flavor-violating ones) are generated. Although the F -component of x contributes to the cosmological constant, it is possible to cancel it out with a fine-tuning of the Kahler potential. If we consider the model given in Ref. [24], for example, vanishing cosmological constant can be realized by a rescaling of the Kahler potential. Gaugino masses can arise from x -dependent gauge kinetic functions (or from direct coupling of z to the gauge kinetic terms as suggested in Ref. [24]). A -parameters are also generated by introducing x -dependent higher-dimensional terms in the superpotential (or by the renormalization-group effects).

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